

# Pentaquark $\Theta^*$ States in the 27-plet from Chiral Soliton Models

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## Abstract

We estimate the mass and the width of pentaquark  $\Theta^*$  states in the 27-plet from chiral soliton models. The calculations show that the mass of  $\Theta^*$  is about 1.60 GeV and the width for the process  $\Theta^* \rightarrow KN$  is less than 43 MeV. We also discuss the search for the existence of  $\Theta^*$  states in physical processes.

*Key words:* pentaquark, chiral soliton model, mass and width, decay modes

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Skyrme's old idea [1] that baryons are solitons has been widely accepted since Witten's topological analysis of the Wess-Zumino term and his clarification in what sense baryons can be considered as classical solitons of effective meson fields [2]. The quantization of the SU(3) Skyrminion not only gives the baryon octet and decuplet, but also predicts new higher baryon multiplets, such as the anti-decuplet, the 27-plet etc [3,4,5]. There are exotic baryon states with strangeness number  $S = +1$  in these higher multiplets, and these states can be interpreted as pentaquark states with minimal five-quark configurations  $uuud\bar{s}$ ,  $uudd\bar{s}$ , and  $uddd\bar{s}$  in the quark language [6]. A number of authors [7,8,9,10] predicted the mass of the lightest pentaquark  $\Theta^+(uudd\bar{s})$  state with  $S=+1$  from chiral soliton models. However, the real boost in searching pentaquark states was due to Diakonov, Petrov and Polyakov's prediction about the mass and the width of  $\Theta^+$  [9]. It seems that recent experiments [11,12,13,14,15] have revealed the existence of  $\Theta^+$  with a

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mass  $M_{\Theta^+} \simeq 1.54$  GeV,  $S = +1$  and a very small width  $\Gamma_{\Theta^+} < 25$  MeV. From the absence of a signal in the corresponding  $pK^+$  invariant mass distribution in  $\gamma p \rightarrow pK^+K^-$  and  $\gamma^* p \rightarrow pK^+K^-$  at the expected strength [14,15], it is suggested that  $\Theta^+$  should be an isoscalar. Thus up to now, experiments have given a surprising support to prediction from chiral soliton models. In Ref. [16], Walliser and Kopeliovich studied the other exotic states in the 27-plet and the  $\overline{35}$ -plet and predicted the mass of the higher  $\Theta$  states in the 27-plet, called  $\Theta^*$  in the following, to be about 1.65/1.69 GeV, provided that the mass of  $\Theta^+$  is at 1.54 GeV. Borisyuk, Faber, and Kobushkin [17] also predicted the mass of  $\Theta^*$  around 1595 MeV and the width at 80 MeV by identifying N(1710) as the member of the anti-decuplet.

In this letter, we calculate both the mass and the width of pentaquark  $\Theta^*$  states from chiral soliton models, and give the predictions based on available experimental observations. The motivation for this is to reveal the existence of pentaquark  $\Theta^*$  states through further experiments, following along the successful prediction from Ref. [9].

Following Ref. [3], the SU(3) symmetric effective action in the large  $N_c$  limit leads to the collective Hamiltonian:

$$\widehat{H} = M_{cl} + \frac{1}{2I_2} \left[ \widehat{C}^{(2)} - \frac{1}{12}(N_c B)^2 \right] + \left( \frac{1}{2I_1} - \frac{1}{2I_2} \right) \widehat{\mathbf{J}}^2, \quad (1)$$

where  $M_{cl}$  is the classical soliton mass;  $\widehat{C}^{(2)} = \sum_{a=1}^8 \widehat{G}_a^2$  is the quadratic (Casimir) operator of the vectorial group SU(3)<sub>v</sub>, in the representation  $(p, q)$ , its eigenvalue  $C^{(2)} = \frac{1}{3}[p^2 + q^2 + pq + 3(p + q)]$ ;  $\widehat{G}_a$  ( $a = 1 - 8$ ) are the generators of SU(3)<sub>v</sub>;  $\widehat{J}_i$  ( $i = 1 - 3$ ) are the generators of the spin group SU(2)<sub>s</sub>;  $I_1$  and  $I_2$  are moments of inertia. Therefore, for the representation  $(p, q)$  of the SU(3)<sub>v</sub> and the spin  $J$ , the eigenvalues of the Hamiltonian are

$$E_J^{(p,q)} = M_{cl} + \frac{1}{6I_2} \left[ p^2 + q^2 + pq + 3(p + q) - \frac{1}{4}(N_c B)^2 \right] + \left( \frac{1}{2I_1} - \frac{1}{2I_2} \right) J(J + 1). \quad (2)$$

From the energy eigenvalues above, it can be argued that the 27-plets with spin 3/2 and 1/2 are the multiplets next to the antidecuplet [5]. The mass differences between the lowest multiplets are

$$\begin{aligned} (1, 1) - (0, 3) : E^{(8)} - E^{(\overline{10})} &= -\frac{3}{2I_2}, \\ (1, 1) - (2, 2) : E^{(8)} - E_{\frac{1}{2}}^{(27)} &= -\frac{5}{2I_2}, \end{aligned}$$

$$\begin{aligned}
(2, 2) - (1, 4) : E_{\frac{3}{2}}^{(27)} - E_{\frac{3}{2}}^{(\overline{35})} &= -\frac{2}{I_2}, \\
(2, 2) - (3, 3) : E_{\frac{3}{2}}^{(27)} - E_{\frac{3}{2}}^{(64)} &= -\frac{7}{2I_2},
\end{aligned}$$

The states of the system will correspond to the baryon states, and wave function  $\Psi_{\nu\nu'}^{(\mu)}$  of baryon  $B$  in the collective coordinates is of the form

$$\Psi_{\nu\nu'}^{(\mu)}(A) = \sqrt{\dim(\mu)} D_{\nu\nu'}^{(\mu)}(A), A \in SU(3), \quad (3)$$

where  $(\mu)$  denotes an irreducible representation of the  $SU(3)$  group;  $\nu$  and  $\nu'$  denote  $(Y, I, I_3)$  and  $(1, J, -J_3)$  quantum numbers collectively;  $Y$  is the hypercharge of  $B$ ;  $I$  and  $I_3$  are the isospin and its third component of  $B$  respectively;  $J_3$  is the third component of spin  $J$ ;  $D_{\nu\nu'}^{(\mu)}(A)$  are representation matrices. However, due to the non-zero strange quark mass, the symmetry breaking Hamiltonian is [18]

$$H' = \alpha D_{88}^{(8)} + \beta Y + \frac{\gamma}{\sqrt{3}} \sum_{i=1}^3 D_{8i}^{(8)} J^i, \quad (4)$$

where the coefficients  $\alpha, \beta, \gamma$  are proportional to the strange quark mass and model dependent, but they are treated model-independently and fixed by experiments in this letter;  $D_{ma}^{(8)}(A)$  is the adjoint representation of the  $SU(3)$  group and defined as:

$$D_{ma}^{(8)}(A) = \frac{1}{2} \text{Tr}(A^\dagger \lambda^m A \lambda^a), \quad (5)$$

and  $\lambda^m$  is the Gell-Mann matrix of the corresponding meson.

Another consequence of the flavor symmetry breaking is that a physical baryon state is no long a pure state belonging to a unique multiplet, but a mixing state with the corresponding members with identical spin and isospin in other multiplets, that is

$$\Psi_{\nu\nu'}(A) = \sum_{\mu} c_{\nu\nu'}^{(\mu)} \Psi_{\nu\nu'}^{(\mu)}(A). \quad (6)$$

From (4), the physical baryon states are of the form by first-order approximation

$$\begin{aligned}
|N\rangle &= |N; 8\rangle + C_{\overline{10}} |N; \overline{10}\rangle + C_{27} |N; 27_{\frac{1}{2}}\rangle, \\
|\Theta^+\rangle &= |\Theta^+; \overline{10}\rangle, \\
|\Theta^*\rangle &= |\Theta^*; 27_{\frac{3}{2}}\rangle + C_{\overline{35}} |\Theta^*; \overline{35}_{\frac{3}{2}}\rangle + C_{64} |\Theta^*; 64_{\frac{3}{2}}\rangle.
\end{aligned}$$

To linear order of  $m_s$ , the coefficients above are given simply by perturbation theory

$$\begin{aligned}
C_{\overline{10}} &= -\frac{1}{3\sqrt{5}}(\alpha + \frac{\gamma}{2})I_2, & C_{27} &= -\frac{\sqrt{6}}{25}(\alpha - \frac{\gamma}{6})I_2. \\
C_{\overline{35}} &= -\frac{3}{4\sqrt{35}}\left(\alpha + \frac{5}{6}\gamma\right)I_2, & C_{64} &= -\frac{3\sqrt{10}}{196}\left(\alpha - \frac{1}{6}\gamma\right)I_2.
\end{aligned} \tag{7}$$

In chiral soliton model, the 27-plet with spin  $\frac{3}{2}$ , lower than that with spin 1/2, is the next multiplet to the anti-decuplet, we only deal with spin-3/2 baryons in this letter, and omit the spin-3/2 index of the notations of particles in the 27-plet as well as energy eigenvalue from now on. The quark content of the exotic pentaquark states are suggested in Fig. 1, and the mass splittings of the isomultiplets in the 27-plet are listed in Table 1.

Table 1. The mass of baryons in the  $\{27\}$  multiplet

Baryon	I	Y	$\langle B H' B\rangle$	Mass (GeV)
exotic pentaquarks				
$\Theta^*$	1	2	$\frac{\alpha}{7} + 2\beta - \frac{5}{14}\gamma$	1.60
$X_{1s}$	2	0	$\frac{5}{56}\alpha - \frac{25}{112}\gamma$	1.68
$X_{2s}$	$\frac{3}{2}$	-1	$-\frac{1}{14}\alpha - \beta + \frac{5}{28}\gamma$	1.87
$\Omega^*$	1	-2	$-\frac{13}{56}\alpha - 2\beta + \frac{65}{112}\gamma$	2.07
$\Delta^*$	$\frac{3}{2}$	1	$\frac{13}{112}\alpha + \beta - \frac{65}{224}\gamma$	1.64
exited states of octet				
$N_{27}$	$\frac{1}{2}$	1	$\frac{1}{28}\alpha + \beta - \frac{5}{56}\gamma$	1.73
$\Sigma_{27}$	1	0	$-\frac{1}{56}\alpha + \frac{5}{112}\gamma$	1.80
$\Xi_{27}$	$\frac{1}{2}$	-1	$-\frac{17}{112}\alpha - \beta + \frac{85}{224}\gamma$	1.96
$\Lambda_{27}$	0	0	$-\frac{1}{14}\alpha + \frac{5}{28}\gamma$	1.86

In experiments, we are interested in the decay  $\Theta^* \rightarrow KN$  which are realized by a pseudoscalar Yukawa coupling. In soliton models, such a coupling can be obtained by Goldberger-Treiman relation, which relates the relevant cou-

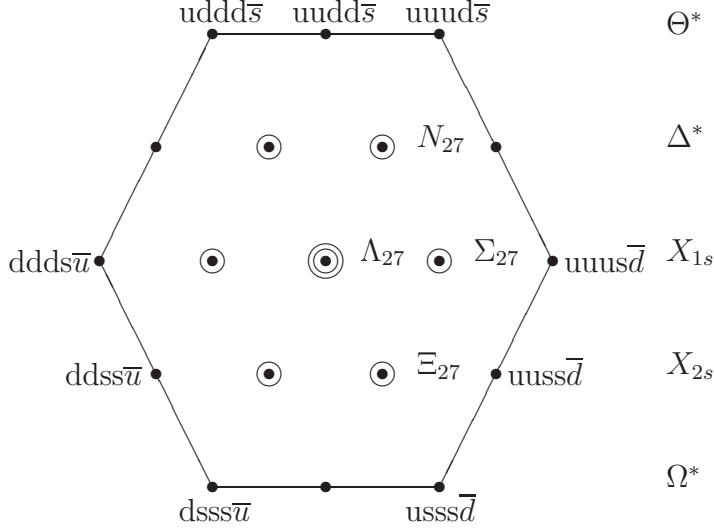


Fig. 1. The quark content of the  $\{27\}$  multiplet baryons.

pling constant to the axial charge [20,21]. And up to  $1/N_c$  order, the coupling operator in the space of the collective coordinates  $A$  has the form [9,21]:

$$\hat{g}_A \propto G_0 D_{m3}^{(8)} - G_1 d_{3ab} D_{ma}^{(8)} J_b - \frac{G_2}{\sqrt{3}} D_{m8}^{(8)} J_3, \quad (8)$$

where  $d_{iab}$  is the  $SU(3)$  symmetric tensor,  $a, b = 4, 5, 6, 7$ , and  $J_a$  are the the generators of the infinitesimal  $SU_R(3)$  rotations.  $G_1, G_2$  are dimensionless constants,  $1/N_c$  suppressed relative to  $G_0$ . Let  $|i\rangle = \Psi_{\rho\rho'}^{(\mu')}(A) + c_i \Psi_{\rho\rho'}^{(\mu_i)}(A)$  denote the state of  $B'$  and  $|f\rangle = \Psi_{\nu\nu'}^{(\mu)}(A) + d_f \Psi_{\nu\nu'}^{(\mu_f)}(A)$  denote the state of  $B$ , where  $c_i$  and  $d_f$  are chosen to be real and are all of  $1/N_c$  order. Then sandwiching  $\hat{g}_A$  between  $|f\rangle$  and  $|i\rangle$  gives the coupling  $g_{BB'm}$  for the decay  $B \rightarrow B'm$

$$g_{BB'm}^2 = \frac{g_0}{(2J_\nu + 1)} \sum_{J_{\nu 3} J_{\rho 3}} \sum_{I_{m3} I_{\rho 3}} \left| \langle f | G_0 D_{m3}^{(8)} - G_1 d_{3ab} D_{ma}^{(8)} J_b - \frac{G_2}{\sqrt{3}} D_{m8}^{(8)} J_3 | i \rangle \right|^2, \quad (9)$$

where we leave  $g_0$  as a constant to be fixed by experiments; and it is averaged over the initial spin( $J_\nu, J_{\nu 3}$ ) and sums over the final spin( $J_\rho, J_{\rho 3}$ ) as well as isospin( $I_m, I_{m3}$  and  $I_\rho, I_{\rho 3}$ ) states. Up to linear order of  $m_s$  and  $1/N_c$  and neglecting the terms proportional to  $m_s/N_c$ , we can rewrite the coupling as

$$g_{BB'm}^2$$

$$\begin{aligned}
&= \frac{g_0 G^2}{3} \times \left\{ \frac{\dim(\mu')}{\dim(\mu)} \left| \sum_{\gamma} \begin{pmatrix} 8 & \mu' & \mu_{\gamma} \\ Y_m I_m & Y_{\rho} I_{\rho} & Y_{\nu} I_{\nu} \end{pmatrix} \begin{pmatrix} 8 & \mu' & \mu_{\gamma} \\ 01 & 1J_{\rho} & 1J_{\nu} \end{pmatrix} \right|^2 + \right. \\
&2c_i \frac{G_0}{G} \frac{\sqrt{\dim(\mu')\dim(\mu_i)}}{\dim(\mu)} \sum_{\gamma} \begin{pmatrix} 8 & \mu' & \mu_{\gamma} \\ Y_m I_m & Y_{\rho} I_{\rho} & Y_{\nu} I_{\nu} \end{pmatrix} \begin{pmatrix} 8 & \mu' & \mu_{\gamma} \\ 01 & 1J_{\rho} & 1J_{\nu} \end{pmatrix} \\
&\times \sum_{\gamma'} \begin{pmatrix} 8 & \mu_i & \mu_{\gamma'} \\ Y_m I_m & Y_{\rho} I_{\rho} & Y_{\nu} I_{\nu} \end{pmatrix} \begin{pmatrix} 8 & \mu_i & \mu_{\gamma'} \\ 01 & 1J_{\rho} & 1J_{\nu} \end{pmatrix} + \\
&2d_f \frac{G_0}{G} \frac{\dim(\mu')}{\sqrt{\dim(\mu)\dim(\mu_f)}} \sum_{\gamma} \begin{pmatrix} 8 & \mu' & \mu_{\gamma} \\ Y_m I_m & Y_{\rho} I_{\rho} & Y_{\nu} I_{\nu} \end{pmatrix} \begin{pmatrix} 8 & \mu' & \mu_{\gamma} \\ 01 & 1J_{\rho} & 1J_{\nu} \end{pmatrix} \\
&\left. \times \sum_{\gamma_i} \begin{pmatrix} 8 & \mu' & \mu_{f\gamma_i} \\ Y_m I_m & Y_{\rho} I_{\rho} & Y_{\nu} I_{\nu} \end{pmatrix} \begin{pmatrix} 8 & \mu' & \mu_{f\gamma_i} \\ 01 & 1J_{\rho} & 1J_{\nu} \end{pmatrix} \right\}, \tag{10}
\end{aligned}$$

where  $G$  can be extracted from  $\langle f | G_0 D_{m3}^{(8)} - G_1 d_{3ab} D_{ma}^{(8)} J_b - \frac{G_2}{\sqrt{3}} D_{m8}^{(8)} J_3 | i \rangle$  [9]. In this approximation, we ignore such cases that some of the SU(3) flavor Clebsch-Gordan coefficients which would multiply  $G_{1,2}$  so that would have  $N_c$  dependence enhancing the naive  $N_c$  power. This is what happens in the case of  $G_{10} = G_0 - \frac{N_c+1}{4} G_1 - \frac{1}{2} G_2$  where the constant  $G_1$ , which is formally  $O(1/N_c)$  with respect to  $G_0$ , is enhanced [23]. The discussion of this effect is beyond the scope of the present paper and an improved width formula is used to discuss the width of the anti-decuplet baryons in Ref. [24]. Then, the width is given by

$$\Gamma(B \rightarrow B' m) = \frac{3g_{BB'm}^2}{4\pi m_B} |\mathbf{p}| \left[ (m_{B'}^2 + \mathbf{p}^2)^{\frac{1}{2}} - m_{B'} \right] \approx \frac{3g_{BB'm}^2}{8\pi m_B m_{B'}} |\mathbf{p}|^3.$$

This formula is the same as that in Refs. [9] and [10] in the non-relativistic case. It is very well satisfied experimentally in the case of the decay of the decuplet baryons, the input data [22] and the numeric results are as listed in Table 2.

Table 2. The best fit at  $g_0=3.84$ 

Decay modes	$m_B, m'_B, m_m$ (MeV)	PDG data	theory values
$\Delta \rightarrow N\pi$	1232, 983.3, 139.6	$\approx 120$	117
$\Sigma^* \rightarrow \Lambda\pi$	1385, 1116, 135	$\approx 34.67$	34.4
$\Sigma^* \rightarrow \Sigma\pi$	1387, 1197, 135	$\approx 4.73$	4.5
$\Xi^* \rightarrow \Xi\pi$	1535, 1321, 135	$\approx 9.9$	10.4

After a trivial calculation, we have the width formulae for  $\Theta$  and  $\Theta^*$  decays:

$$\Gamma(\Theta^+ \rightarrow KN) = g_0 \frac{(G_0 - G_1 - \frac{1}{2}G_2)^2}{40\pi m_{\Theta^+} m_N} |\mathbf{p}|^3 [1 + \frac{G_0}{G_0 - G_1 - \frac{1}{2}G_2} (\frac{\sqrt{5}}{2} C_{\overline{10}} - \frac{7\sqrt{6}}{12} C_{27})], \quad (11)$$

$$\Gamma(\Theta^* \rightarrow KN) = g_0 \frac{(G_0 - \frac{1}{2}G_1)^2}{54\pi m_{\Theta^*} m_N} |\mathbf{p}|^3 [1 - \frac{G_0}{G_0 - \frac{1}{2}G_1} (\frac{\sqrt{5}}{2} C_{\overline{10}} + \frac{3\sqrt{6}}{28} C_{27})]. \quad (12)$$

Ref. [25] first reported evidence for the existence of a narrow  $\Xi^-\pi^-$  baryon resonance with mass of  $1.862 \pm 0.003$  GeV and width below the detector resolution of about 0.018 GeV, and this state is considered as a candidate for the pentaquark  $\Xi_{\frac{3}{2}}^{--}$ . If we take both  $\Theta^+$  and the candidate for  $\Xi_{3/2}$  [25] as members of the anti-decuplet and solve the following equations

$$\left\{ \begin{array}{l} \frac{1}{I_1} = \frac{2}{3} [E^{(10)} - E^{(8)}] = \frac{2}{3} [m_{\Sigma^*} - \frac{1}{2}(m_{\Lambda} + m_{\Sigma})] = 154 \text{ MeV}; \\ \alpha + \frac{3}{2}\gamma = 5(m_{\Lambda} - m_{\Sigma}) = -385 \text{ MeV}; \\ \frac{1}{8}\alpha + \beta - \frac{5}{16}\gamma = m_{\Delta} - m_{\Sigma^*} = -153 \text{ MeV}; \\ E^{(\overline{10})} + (\frac{1}{4}\alpha + 2\beta - \frac{1}{8}\gamma) = m_{\Theta^+} = 1540 \text{ MeV}; \\ E^{(\overline{10})} - (\frac{1}{8}\alpha - \beta + \frac{1}{16}\gamma) = m_{\Xi_{3/2}} = 1860 \text{ MeV}; \\ E^{(\overline{10})} - \frac{3}{2I_2} = E^{(10)} - \frac{3}{2I_1} = 1154.5 \text{ MeV}; \\ E^{(27)} - E^{(\overline{10})} + \frac{1}{2I_2} = \frac{3}{2I_1} = 230.5 \text{ MeV}; \end{array} \right. \quad (13)$$

we get

$$\begin{aligned}
m_{\Theta^*} &= 1.60 \text{ GeV}; \quad E^{(27)} = 1.785 \text{ GeV}; \quad 1/I_2 = 399 \text{ MeV}; \\
\alpha &= -663 \text{ MeV}; \quad \beta = -12 \text{ MeV}; \quad \gamma = 185 \text{ MeV}; \\
C_{\overline{10}} &= 0.21; \quad C_{27} = 0.17; \quad C_{\overline{35}} = 0.16; \quad C_{35} = 0.08.
\end{aligned}$$

We find that the values of  $C_{\overline{10}}$  and  $C_{27}$  are the same as those in Ref. [26]. The masses of baryons for the 27-plet in the parameters above are listed in Table 1. Using the results above, we can calculate the width for  $\Theta^* \rightarrow KN$ :

$$\Gamma(\Theta^* \rightarrow KN) = g_0 \frac{(G_0 - \frac{1}{2}G_1)^2}{54\pi m_{\Theta^*}^* m_N} |\mathbf{p}|^3 \left[ 1 - \frac{G_0}{G_0 - \frac{1}{2}G_1} \left( \frac{\sqrt{5}}{2} C_{\overline{10}} + \frac{3\sqrt{6}}{28} C_{27} \right) \right] \leq 43 \text{ MeV}. \quad (14)$$

Thus, soliton model gives a stringent restriction on the width of  $\Theta^*$  for the process  $\Theta^* \rightarrow KN$ , and if, as reported by recent experiments,  $\Gamma_{\Theta^+} < 25 \text{ MeV}$ , the width for  $\Theta^* \rightarrow KN$  will be less than 43 MeV. The predicted width could be more narrow if a smaller input width  $\Gamma_{\Theta^+}$  is used.

In the pentaquark  $\Theta^*$  triplet,  $\Theta^{*++}$  and  $\Theta^{*+}$  may be easily measured. The search for  $\Theta^{*+}$  is similar to  $\Theta^+$  through the decay modes  $\Theta^{*+} \rightarrow K^+ n$  and  $K^0 p$  with approximately same magnitudes [27]. There have been suggestions for search of pentaquark  $\Theta^{++}(uuuds\bar{s})$  state in virtual and real photon scattering on the proton target [6,28]. Provided with the ranges of the predicted mass and width, the existence of  $\Theta^{*++}$  may be revealed through the decay mode  $\Theta^{*++} \rightarrow K^+ p$  in various processes, such as:

Photon-nucleon collisions

$$\gamma p \rightarrow \Theta^{*++} K^-; \quad \Theta^{*++} \rightarrow p K^+;$$

Nucleon-nucleon collisions

$$pp \rightarrow p K^- \Theta^{*++}; \quad \Theta^{*++} \rightarrow p K^+;$$

Pion-nucleon collisions

$$\pi^+ p \rightarrow \overline{K}^0 \Theta^{*++}; \quad \Theta^{*++} \rightarrow p K^+;$$

Electron(virtual photon)-nucleon collisions

$$ep \rightarrow e' K^- \Theta^{*++}; \quad \Theta^{*++} \rightarrow p K^+.$$

$\Theta^{*0}$  may be revealed through the decay mode  $\Theta^{*0} \rightarrow K^0 n$  in the above corresponding processes with the target changed from proton to neutron. The correlations between the constructed  $KN$  invariant masses of  $\Theta^{*++}$ ,  $\Theta^{*+}$ , and  $\Theta^{*0}$  decays can test whether the corresponding states are belong to the pentaquark  $\Theta^*$  triplet suggested in this work.

In summary, calculations from the chiral soliton model show that, the pentaquark  $\Theta^*$  states in the 27-plet with spin 3/2, have a mass around 1.60 GeV



and a width for  $\Theta^* \rightarrow KN$  less than 43 MeV. The existence of these pentaquark  $\Theta^*$  states can be revealed by the decay modes  $\Theta^* \rightarrow KN$  in various physical processes.

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